

# A Quirky Probe of Neutral Naturalness

Zackaria Chacko,<sup>\*</sup> David Curtin,<sup>†</sup> and Christopher B. Verhaaren<sup>‡</sup>  
*Maryland Center for Fundamental Physics, Department of Physics,*  
*University of Maryland, College Park, MD 20742-4111 USA*  
 (Dated: December 21, 2015)

We consider the signals arising from top partner pair production at the LHC as a probe of theories of Neutral Naturalness. We focus on scenarios in which the top partners carry electroweak charges, such as Folded SUSY or the Quirky Little Higgs. In this class of theories the top partners are charged under a mirror color group whose lightest states are hidden glueballs. The top partners, once produced, form a quirky bound state that de-excites promptly and annihilates into hidden glueballs. These glueballs can decay through mixing with the Higgs, giving rise to striking displaced signatures at the LHC. We show that the displaced signals arising from top partner pair production constitute the primary discovery channel for this class of theories in most regions of parameter space.

The Standard Model (SM) of particle physics is a theoretical triumph whose last missing ingredient, a Higgs boson with approximately the expected couplings, was discovered by both experiments at the Large Hadron Collider (LHC) in 2012 [1, 2]. Despite its experimental successes, it suffers from the well-known *hierarchy problem* [3]: there are large, quadratically divergent quantum corrections to the Higgs mass parameter which have to cancel against the bare mass term to obtain the physically observed mass of 125 GeV. From a Wilsonian Effective Field Theory point of view, this leads us to expect new degrees of freedom that couple to the Higgs and cancel the SM loops. Since the largest quadratic divergence arises from the top loop, we expect new physics that cancels this at a scale below  $\sim$  TeV. Otherwise, the theory is *tuned* or *unnatural*.

In theories such as supersymmetry [4] or the Little Higgs [5–8], the top loop is canceled by *top partner states* that are related to the top quark by a symmetry transformation. The symmetry relates the coupling of the top partner to the Higgs to that of the top quark, thereby enforcing the cancellation. These top partners carry SM color charge, leading to copious production at the LHC if their masses lie below the TeV scale. While the absence such a discovery at the first run of the LHC can be explained by kinematic blind-spots or non-minimal scenarios [9–19], these null results lead to some tension with naturalness.

In theories of *Neutral Naturalness* [20–22] the top loop is canceled by introducing top partners that do not carry SM color charge. This can occur when the symmetry that protects the Higgs mass does not commute with SM color. This class of theories is naturally consistent with the LHC limits on colored particles. They also offer a more general framework in which to consider the robust experimental consequences of naturalness. The phenomenology of these models can be very rich, and is, in general, radically different from scenarios involving colored top partners.

In most realistic theories of Neutral Naturalness the top partners are charged under a mirror copy of QCD. The top partners can carry SM electroweak (EW) charge, as in the case of Folded Supersymmetry (FSUSY) [21] and the Quirky Little Higgs (QLH) [22], or remain SM singlets, as in the Twin

Higgs (TH) [20, 23, 24] family of theories. These models have rich implications for cosmology [25–31], and possibly also for flavor [32]. UV completions [33–41] are required at scales of order 5 – 10 TeV to protect against higher loop effects. At these energies the full protection mechanism of the theory is expected to become apparent. This strongly motivates the construction of future lepton and 100 TeV colliders [42, 43].

At the LHC, the most promising signals of Neutral Naturalness are *displaced signatures* that arise when these theories realize a specific Hidden Valley [44–47] scenario. This was first explicitly pointed out in [26] in the context of the Fraternal Twin Higgs (FTH) model. If there is no light matter charged under the mirror color group, the lightest hadrons in the hidden QCD spectrum are glueballs [48]. Mirror gluons couple to the Higgs via a dimension-6 operator generated by the top partner loop [49], in direct analogy to SM tops and gluons. This operator generates mixing between the  $0^{++}$  glueballs and the Higgs, allowing these states to decay to SM particles, primarily  $b\bar{b}$  and  $\tau\tau$ , through the Higgs portal. These decays are slow on collider timescales, involving characteristic decay lengths of  $\mu\text{m}$  - km, which can be reconstructed in the LHC detectors. This implies that any process that results in the production of these mirror glueballs, such as decays of the Higgs to a pair of mirror gluons, will lead to striking signals at the LHC.

This signature is particularly well-motivated in the case of theories such as FSUSY and QLH, in which there is no light matter charged under the hidden color group. In TH models this signature can be realized, but only if the light mirror quarks are removed from the low-energy spectrum, as is the case for the FTH model. In these theories, renormalization group arguments motivate the lightest glueball mass to be in the 10 - 60 GeV range, allowing for exotic Higgs decays. Displaced searches at the LHC for mirror glueballs arising from Higgs decays are then projected to be sensitive to 600 - 800 GeV top partners at the end of run 2, and TeV-scale top partners by the end of the HL-LHC [50], as shown in Fig. 1. Even the first 20 fb<sup>-1</sup> of 13 TeV data offer a reach of a few hundred GeV [51]. By comparison, precision measurements of  $h \rightarrow \gamma\gamma$  will only be sensitive to top partner masses of a few

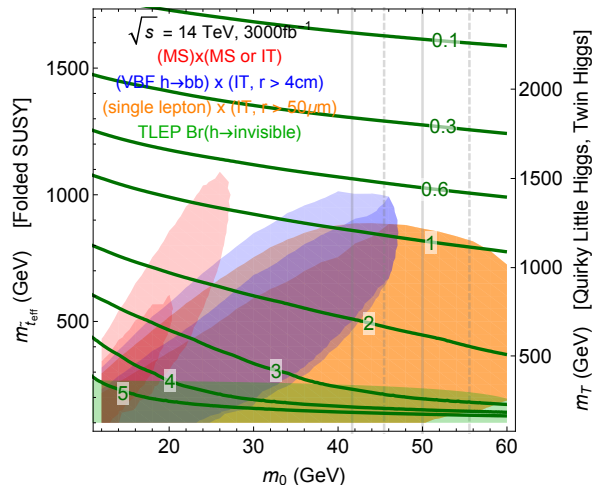


FIG. 1. *Shaded regions*: Projected sensitivity of different proposed displaced vertex searches, at the 14 TeV LHC with  $3000 \text{ fb}^{-1}$ , to long-lived mirror glueballs from exotic Higgs decays in theories of Neutral Naturalness [50]. The bounds are expressed as a function of the lightest glueball mass (horizontal axis) and top partner mass  $m_{t_{\text{eff}}}$  in FSUSY for degenerate unmixed stops (left axis) and  $m_T$  in TH/QLH (right axis). Light shading represents the factor of  $\sim 10$  uncertainty in the number of the lightest  $0^{++}$  glueballs produced during mirror hadronization when heavier glueball states are kinematically accessible. Higgs decay is assumed to be two-body, which underestimates sensitivity, especially for low  $m_0$ . See [50] for details. *Green contours*: Maximally conservative estimate of the number of glueballs produced from top partner pair production and annihilation in the Quirky Little Higgs model, normalized to the rate from exotic Higgs decays. See Eq. (5).

100 GeV [52]. This serves to illustrate the exquisite sensitivity of exotic Higgs decays to new physics [53].

In this letter we investigate another promising avenue for probing theories of Neutral Naturalness: the displaced signatures arising from the mirror glueballs that result from *direct top partner production*. In theories such as FSUSY and QLH, top partners can be pair produced through Drell-Yan (DY) or vector boson fusion (VBF) processes with sizable rate at the LHC. The top partner pair forms a so-called *quirky bound state* [54–56], since the mirror gluon string connecting them cannot snap by exciting light quark pairs out of the vacuum. The bound state promptly de-excites via soft emission of glueballs and photons [56–61] and can annihilate to a pair of mirror gluon jets. Top partner pair production can therefore result in mirror gluon dijet production. The glueballs resulting from mirror hadronization can then give rise to events with multiple displaced vertices, which are even more conspicuous than the events arising from exotic Higgs decays.

Quirky signals of FSUSY have been previously considered in [57]. However, that study focused on pair production of first and second generation partners and annihilation into  $W\gamma$ . The mass of those states is not as closely connected to naturalness as that of the top partners, and this final state has much more

SM background than displaced decays.

In this paper we show that pair production of top partners, with subsequent annihilation into mirror glueballs, can constitute the discovery signature of Neutral Naturalness at the LHC in large regions of parameter space. A key challenge in making this prediction is the quantitative treatment of mirror hadronization, the details of which are not well understood for a pure  $SU(3)$  gauge theory. Even so, we demonstrate how to consistently parameterize our ignorance about the non-perturbative physics of the hidden sector, allowing for a systematic study of the signatures. We identify the regions of parameter space in which direct production is definitely superior to exotic Higgs decays as a probe of top partner mass, even with maximally pessimistic assumptions about the hadronization of the mirror gluon jets. A full exploration of the signature space, which can include final states with many  $b\bar{b}$  pairs, displaced vertices, and missing energy, and which might allow for the measurement of top partner masses and couplings, will be explored in a detailed follow-up publication [62].

It is also interesting to consider top partner pair production in FTH models. Due to their SM singlet nature, this process occurs via an off-shell Higgs only, with significantly smaller cross section than for EW-charged partners. Furthermore, the mirror partners of the  $b$ -quark are present in the low-energy spectrum, leading the top partner pair to (mirror-) beta-decay to a mirror-bottomonium state which then decays to the SM either directly or via mirror glueballs. In either case, the number and energy of the produced glueballs is much lower than for EW partners. These two factors likely mean that a 100 TeV collider is required to study this signature, which would serve as a powerful diagnostic of the hidden sector and be complementary to direct production of multi-TeV states in the UV completion [43]. That being said, the presence of mirror bottomonium states in the FTH model makes the exotic Higgs decay phenomenology much more varied and rich than for FSUSY and QLH [26]. We will quantitatively study the LHC reach for these signatures beyond the glueball case explored by [50] in a future study [63].

We now analyze top partner pair production in FSUSY and the QLH. While these models serve as useful theory benchmarks, our conclusions are general and should apply to any scenario with EW-charged top partners charged under a mirror QCD force.

**FOLDED SUPERSYMMETRY** — In the 5D FSUSY theory [21], all QCD-charged fields of the MSSM, and the  $SU(3)_c$  gauge sector itself, are duplicated into two sectors  $A$  (SM) and  $B$  (mirror) at some multi-TeV scale, with couplings related by a discrete  $\mathbb{Z}_2$  symmetry. At energies  $\lesssim$  TeV, the electroweak and Higgs sectors are similar to the 4D MSSM with decoupled gauginos. However, only our  $A$ -sector quarks and  $B$ -sector mirror squarks have light zero modes. This realizes an accidental low-energy SUSY limit, in which the quadratically divergent top contribution to the Higgs mass is cancelled by the mirror-sector stops, which carry the same Yukawa couplings and electroweak gauge quantum numbers as conventional stops, but are charged under the mirror QCD

force.

For the purpose of this phenomenological discussion, the expressions for the lightest squark masses in FSUSY [21] can therefore be taken to be the same as in the MSSM [4]. The light mirror hadrons are glueballs, as described above. Following the methodology of [50], we concentrate on the displaced signatures of the lightest  $0^{++}$  glueball.

The minimal FSUSY model predicts a slightly lower soft mass for the right-handed stop  $\tilde{t}_R$  than for the left-handed stop  $\tilde{t}_L$ . There are no tree-level  $A$ -terms, but the stop mixing angle is difficult to predict since it depends on the  $\mu$ -parameter in the MSSM-like Higgs sector. The lack of viable electroweak symmetry breaking [64] in the minimal model necessitates additional soft mass contributions, which are likely to arise from brane-localized Kahler operators [65–67]. Even with additional soft mass terms, the physical Higgs mass in FSUSY is expected to be smaller than the observed value for stops lighter than a TeV, much as in the MSSM. Additional Higgs mass contributions can be generated by, for example, non-decoupling  $D$ -terms as shown in [52], or by adding extra singlet scalars as in the NMSSM. The details of how the physical Higgs mass is generated has no significant effect on our discussion of the quirk signature. We therefore regard the stop masses and mixing angle  $\theta_t$  as free parameters. We assume the lightest squark to be a stop, as motivated by naturalness, but if this is not the case then a similar discussion applies to the quirky signatures of the lightest squark state, with the exception that the likely near-degenerate nature of first and second generation squarks implies the additional  $W\gamma$  signal discussed in [57].

We now outline how to estimate the signal from top partner pair production. The stops are produced electroweakly, with a cross section that can be readily computed, e.g. in MadGraph [68]. For stop masses below (above)  $\sim \text{TeV}$ , DY (VBF) type production dominates. The stops form a quirky bound state, connected by a mirror gluon flux tube that cannot break due to the absence of light mirror QCD-charged matter. The bound state sheds energy by emission of soft glueballs and photons, which randomizes its orbital angular momentum in a random walk. Annihilation at high orbital angular momentum is strongly suppressed [56], favoring annihilation after the state is almost fully de-excited, with the non-relativistic stops forming  $s$ -wave stoponium  $\eta_{\tilde{t}}$  before annihilating. The annihilation branching fractions can be easily computed by adapting the results of [10]. The mirror di-gluon final state usually dominates with a branching ratio of  $\sim 50 - 80\%$  in most of our parameter space of interest. This is due to the large hidden sector QCD coupling and gluon multiplicity. For large stop mass splittings and mixings, however, annihilation into two 125 GeV Higgs bosons can dominate (see also [69]), while  $WW$ ,  $ZZ$  are produced  $\sim 10\%$  of the time. These SM final states may be particularly useful for precise mass measurements. Details will be included in [62]. Here we focus on the mirror gluon final state due to the low background of displaced searches.

If lighter states (like the sbottom) are available, one or both

of the stops may  $\beta$ -decay, which adds leptons to the mirror gluon jet signature. Whether  $\beta$  decay occurs before annihilation depends on the mass-splitting [57, 59], but here we concentrate on the case where the lightest stop is pair produced and cannot  $\beta$  decay. We will indicate the regions of parameter space where this is not the case.

As part of our conservative estimate, we will ignore the soft emission of photons and glueballs during de-excitation as part of the signal, and concentrate solely on the mirror gluon jet-pair that is created when the quirk state annihilates.

**Mirror Gluon Jets** — We now discuss the properties of mirror gluon jets in detail. The perturbative shower evolution of the mirror gluons proceeds very similarly to the SM, except without quarks and with a coupling  $\alpha_s^B$  that is a modest  $\mathcal{O}(1)$  factor higher than the SM  $\alpha_s^A$  due to differences in RG evolution [50]. The largest possible opening angle of the first gluon splitting in the shower,

$$\frac{2m_J^2}{E_{\text{gluon}}^2} \approx 1 - \cos \theta_{\text{max}} \approx \frac{1}{2} \theta_{\text{max}}^2, \quad (1)$$

determines the rough “width” of the jet. The jet mass  $m_J$  distribution will be similar to the SM in  $e^+e^-$  collisions, which has been the subject of many perturbative studies, e.g. [70, 71]. The average mass of the heavier jet scales like the center of mass energy  $E_{\text{CM}}$

$$\langle m_J^2 \rangle \simeq \frac{\alpha_s^B(E_{\text{CM}}^2)}{\pi} E_{\text{CM}}^2. \quad (2)$$

Since  $E_{\text{CM}} \gtrsim 2m_{\tilde{t}}$  in quirky annihilation, we expect the overall width of these jets to be similar to high-energy SM jets. For the purpose of this discussion, we are therefore justified in treating these mirror di-jets as pencil-like. We do not attempt to compute the precise jet width, which relies on incalculable details of the hadronization and has a highly subdominant effect on our signal acceptance compared to the other uncertainties.

Having settled the issue of rough jet shape, the most important question is that of hadron multiplicity: we need to know how many mirror glueballs are produced in each jet (which determines glueball momentum), and what fraction is the lightest  $0^{++}$  that gives rise to displaced vertices. Unfortunately, the details of pure gauge hadronization are completely unknown, and not reliably calculable by any known method. We therefore adopt a way of parameterizing our ignorance that allows us to systematically consider the range of possibilities for hadronization, in a way that is parametrically transparent and accurate with  $\mathcal{O}(1)$  precision for the overall signal estimate, while factorizing from the “hard” theory parameters like top partner mass and glueball mass.

Glueball multiplicities are encoded in the nonperturbative fragmentation function for the mirror gluon. While its magnitude is unknown, we can perturbatively compute how it changes with scale using the DGLAP evolution equation [72]. The multiplicity of any given hadron (in the massless limit)

scales as

$$\langle n(E_{\text{CM}}^2) \rangle \propto \exp \left( \frac{12\pi}{33} \sqrt{\frac{6}{\pi \alpha_s^B(E_{\text{CM}}^2)}} + \frac{1}{4} \ln \alpha_s^B(E_{\text{CM}}^2) \right), \quad (3)$$

where  $\alpha_s^B(E_{\text{CM}})$  is determined by the glueball mass (which fixes  $\Lambda_{\text{QCD}}^B$ ) and the assumption that the stop is the lightest particle charged under mirror QCD.

This allows us to define two quantities,  $N_G(E_{\text{CM}})$  and  $r_{G_0}$ .  $N_G(E_{\text{CM}})$  is the *total* number of glueballs produced, on average, by the hadronization of a mirror gluon. Its dependence on the center-of-mass energy  $E_{\text{CM}} \gtrsim 2m_{\tilde{t}}$  is given by Eq. (3), and fixed for all events and different stop masses once  $N_G$  has been specified at a given  $E_{\text{CM}}$ .  $r_{G_0}$  is the fraction of those glueballs that are in the lightest  $G_0 = 0^{++}$  state. In the approximation that hadron masses are much smaller than  $E_{\text{CM}}$  (which is appropriate for our analysis) we expect that  $r_{G_0}$  is approximately constant for different  $E_{\text{CM}}$ .

We can therefore parameterize our ignorance about mirror hadronization by considering the 2D-parameter space of possible values for  $(N_G^0, r_{G_0})$ , where  $N_G^0 = N_G(E_{\text{CM}})$  for some fixed  $E_{\text{CM}}$ . This space is bounded:  $N_G^0 \geq 1$  but smaller (per degree of freedom) than, say, charged hadron production in the SM, since glueballs are heavier and more expensive to produce. (There is also an upper bound for light stops due to the non-negligible mass of mirror glueballs.) Similarly,  $r_{G_0} \leq 1$  and likely larger than, say, 0.1, since  $0^{++}$  is the lightest of the  $\sim 12$  stable glueball state. In fact, [73] estimates the relative abundances of the different glueball species using thermal partition functions at  $T \sim \Lambda_{\text{QCD}}^B$ , which yields  $r_{G_0} \sim 0.5 - 0.6$ , with most of the remaining glueballs being  $2^{++}$  and highly subdominant fraction of the other states. While this result is unlikely to be correct in detail, it reinforces our intuition that  $r_{G_0}$  cannot be very small.

**Signal Estimate** — It is now a straightforward exercise to estimate the number of  $0^{++}$  glueballs produced in each top partner pair production event. In [62] we will use this formalism to explore the landscape of possible quirk signals in detail. Here we merely motivate that study by comparing the number of produced glueballs in top partner pair production to exotic Higgs decays, as discussed in [50].

We assume, very pessimistically, that the number of glueballs produced in the annihilation of two 200 GeV stops is the same as the number of glueballs produced in the decay of the 125 GeV Higgs boson. Setting  $E_{\text{CM}} = 2m_{\tilde{t}}$ , this sets  $N_G^0 = N_G(400 \text{ GeV}) = 1$ , while  $N_G \approx 2.5$  for  $m_{\tilde{t}_1} = 2 \text{ TeV}$ . In computing the ratio of produced number of  $0^{++}$  glueballs in the two processes,  $r_{G_0}$  is about the same and drops out. We can therefore compare the signal rates by computing the ratio

$$R_{\text{FSUSY}} = \frac{\sigma_{\text{DY+VBF}}(pp \rightarrow \tilde{t}_1 \tilde{t}_1) \text{Br}(\eta_{\tilde{t}} \rightarrow g_B g_B) N_G(2m_{\tilde{t}_1})}{\sigma_{\text{VBF}}(pp \rightarrow h) \varepsilon_{\text{VBF}} \text{Br}(h \rightarrow g_B g_B)} \quad (4)$$

where we have assumed VBF Higgs production, and  $\varepsilon_{\text{VBF}} \approx 20\%$  is a generous estimate of the acceptance for VBF triggers

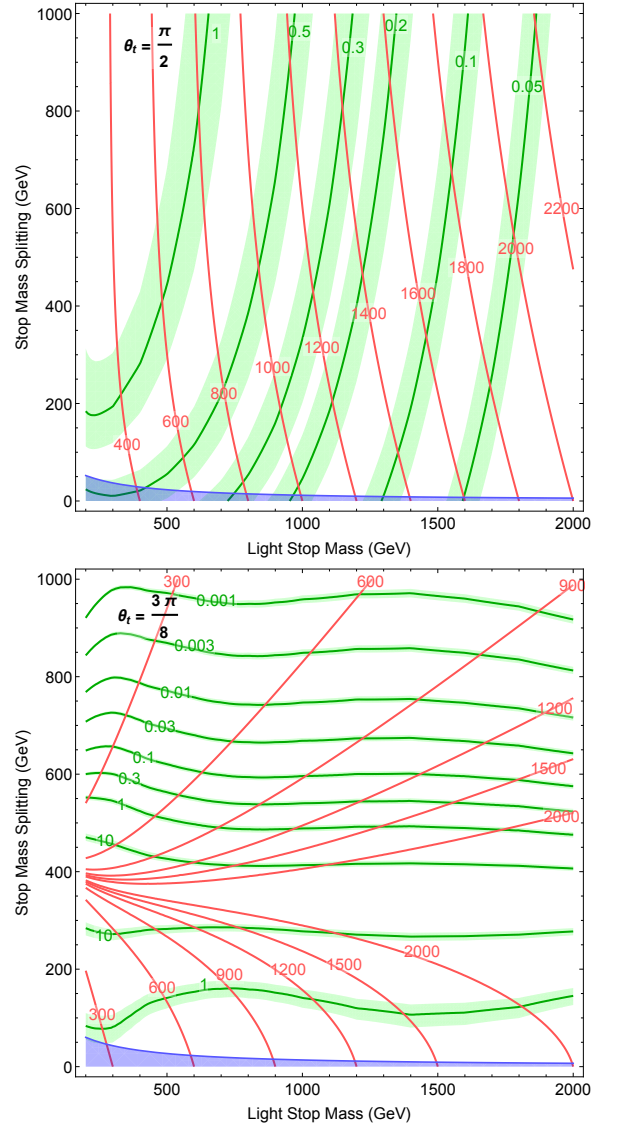


FIG. 2. Green contours show  $R_{\text{FSUSY}}$ , conservatively estimating the number of glueballs produced in stop pair production normalized to exotic Higgs decays in FSUSY, see Eq. (4), as a function of lightest stop  $\tilde{t}_1$  mass and mass splitting, for purely RH light stop (top) and some mixing (bottom). Green shading shows the effect of varying the glueball mass  $m_0$  from 15 GeV (right edge of band) to 50 GeV (left edge). Red contours show  $m_{\text{eff}}$ , which corresponds to the left vertical axis in Fig. 1. Blue shading indicates where  $b_L$  is lighter than  $\tilde{t}_1$ , allowing for the possibility of  $\beta$  decay.

[50]. The pair production cross section is computed at parton-level in MadGraph [68]. Multiplication of the signal rate by the glueball multiplicity per jet is justified for this simple estimate, since reconstruction efficiencies for displaced vertices in the tracker are in the  $\mathcal{O}(10\%)$  range [50], and we assume the number of glueballs per mirror gluon jet to be small. This assumption also implies that each glueball has enough momentum for its decay products to pass various jet or jet + MET trigger thresholds, allowing us to omit a trigger efficiency for

the signal in the numerator. (In the case of many glueballs per jet, trigger acceptance may be slightly reduced but this will be more than compensated by the larger number of events in which displaced vertices can actually be reconstructed.)

This ratio is shown as the green contours in Fig. 2 for two stop mixing angles. There are large regions where this ratio is larger than 1, indicating more displaced vertices from top partner pair production than exotic Higgs decays. In fact, given how conservative our estimate of  $N_G^0$  is, pair production is likely to be the superior discovery channel even when  $R_{\text{FSUSY}}$  is somewhat smaller than 1. To understand the gain in top partner mass reach, we also show contours of  $m_{\text{eff}}$  (red), which corresponds to the left vertical axis of Fig. 1. The bounds on  $m_{\text{eff}}$  from exotic Higgs decays are  $\sim 1$  TeV at the HL-LHC, and a factor of 10 in signal corresponds to  $\sim 200$  GeV in reach. For unmixed RH stops (top plot), pair production is the discovery channel for masses  $< 500 - 1000$  GeV. Pair production is even more important for mixed stops (bottom plot), where exotic Higgs decays are suppressed by cancellations. In fact, for the moderately mixed example shown, quirky pair production is competitive or dominant for all  $m_{\tilde{t}_1} < 2$  TeV and mass splitting below  $\sim 500$  GeV. Note, however, that the annihilation branching fraction to mirror gluons becomes small for large mass splittings. In that region of parameter space, di-higgs searches may have greater sensitivity. In the case of purely LH stops the quirk state is likely to  $\beta$ -decay to mirror-sbottoms, so the resulting leptons increase the conspicuousness of the signal.

Our results are robust, even when corrections arising from the physics that raises the Higgs mass to 125 GeV are taken into account. For example, additional non-decoupling  $D$ -terms would lead to new quartic couplings of the form  $|\tilde{t}|^2|H|^2$ , and also add a multi-TeV scale  $Z'$ . The changed coupling might modify the quirk annihilation branching fractions to Higgs bosons by at most an  $\mathcal{O}(1)$  factor, but only when this branching fraction is not generated by large  $A$ -terms and hence very subdominant to annihilation into mirror gluons. The additional  $Z'$  has to have multi-TeV mass and can be neglected for quirk production. On the other hand, additional singlets as in the NMSSM would change the various couplings to the Higgs by a small mixing angle, which again does not significantly change our conclusions.

**QUIRKY LITTLE HIGGS** — The Quirky Little Higgs model [22] shares features of FSUSY and TH. It features a vector-like fermion top partner, which is an  $SU(2)_L$  singlet carrying hypercharge  $2/3$ . Signal can be estimated in an exactly analogous fashion to the FSUSY discussion above, with a few modifications. There is no mass splitting between different top partner states, allowing us to plot results in the same  $(m_0, m_T)$  plane as the exotic Higgs decay bounds of [50]. VBF production is not competitive with DY and can be omitted. A slight complication is that the quirk can now annihilate as either a spin singlet  $^1S_0$  (which can annihilate to di-gluons) or triplet  $^3S_1$  (which annihilates to at least three gluons). The relevant quirk annihilation widths can be adapted from [74] while accounting for the fact that the top partners do not re-

ceive most of the mass from the light Higgs vev and do not couple axially to the  $Z$ -boson. The singlet width is about 100 times larger than the triplet width, giving two mirror gluons as the highly dominant decay channel. We then compute the analogous ratio to Eq. (4):

$$R_{\text{QLH}} = \frac{\sigma_{\text{DY}}(pp \rightarrow T\bar{T})\text{Br}(^1S_0, ^3S_1 \rightarrow g_B g_B)N_G(2m_T)}{\sigma_{\text{VBF}}(pp \rightarrow h)\varepsilon_{\text{VBF}}\text{Br}(h \rightarrow g_B g_B)} \quad (5)$$

and show it as green contours overlaid on the HL-LHC projected Exotic Higgs decay bounds in Fig. 1. Note that top partner mass in the QLH model is on the right vertical axis of that plot. Even with these pessimistic assumptions for the number of glueballs produced in top partner annihilation, pair production yields more signal events than exotic Higgs decays in the entire region of parameter space where the latter have sensitivity. Therefore, top partner pair production is the main discovery channel for Neutral Naturalness in the Quirky Little Higgs scenario at the LHC.

**CONCLUSIONS** — In this letter we analyzed top partner pair production in theories of Neutral Naturalness at the LHC. This is particularly motivated when the top partners have EW charge, as in Folded SUSY or the Quirky Little Higgs. Mirror glueballs are at the bottom of the mirror spectrum, and the top partners form a quirky bound state which annihilates into jets of mirror gluons. The details of hadronization in a pure  $SU(3)$  gauge theory are unknown, but we devise a simple way of parameterizing our ignorance of mirror hadronization in a way that is transparent, allows for signal estimates at the  $\mathcal{O}(1)$  level, can be applied consistently event-by-event, and factorizes from perturbative theory parameters like the top partner mass.

Our analysis shows that production of mirror glueballs in top partner pair production, which gives rise to displaced vertex signals at the LHC, can be competitive or dominant to glueball production in exotic Higgs decays as analyzed in [50, 51]. This means that top partner pair production is the likely discovery channel of Neutral Naturalness in the Quirky Little Higgs model, and in many FSUSY scenarios, in particular when the reach of exotic Higgs decays is degraded due to cancellations arising from stop mixing.

Top partner pair production can give rise to a rich landscape of signatures, including many displaced vertices (a particular realization of the Emerging Jets scenario [75]), missing energy, and  $b\bar{b}$  pairs. It also offers the tantalizing possibility of measuring the top partner masses and couplings directly, which could confirm the Neutral Naturalness mechanism for solving the Little Hierarchy Problem. We will outline the detailed exploration of these signals in a follow-up publication [62].

**Acknowledgements:** We thank Markus Luty, Yuhsin Tsai, George Sterman, and the participants of the 2015 CERN-CKC Neutral Naturalness workshop for useful discussion. We thank Olivier Mattelaer for helping us with MadGraph. Z.C., D.C., and C.V. are supported by National Science Foundation grant No. PHY-1315155 and the Maryland Center for Funda-

mental Physics.

\* zchacko@umd.edu

† dcurtin1@umd.edu

‡ cver@umd.edu

- [1] G. Aad *et al.* (ATLAS), Phys.Lett. **B716**, 1 (2012), arXiv:1207.7214 [hep-ex].
- [2] S. Chatrchyan *et al.* (CMS), Phys.Lett. **B716**, 30 (2012), arXiv:1207.7235 [hep-ex].
- [3] V. F. Weisskopf, Phys. Rev. **56**, 72 (1939).
- [4] S. P. Martin, Adv.Ser.Direct.High Energy Phys. **21**, 1 (2010), arXiv:hep-ph/9709356 [hep-ph].
- [5] N. Arkani-Hamed, A. G. Cohen, and H. Georgi, Phys.Lett. **B513**, 232 (2001), arXiv:hep-ph/0105239 [hep-ph].
- [6] N. Arkani-Hamed, A. Cohen, E. Katz, A. Nelson, T. Gregoire, *et al.*, JHEP **0208**, 021 (2002), arXiv:hep-ph/0206020 [hep-ph].
- [7] N. Arkani-Hamed, A. Cohen, E. Katz, and A. Nelson, JHEP **0207**, 034 (2002), arXiv:hep-ph/0206021 [hep-ph].
- [8] M. Schmaltz, JHEP **0408**, 056 (2004), arXiv:hep-ph/0407143 [hep-ph].
- [9] S. P. Martin, Phys.Rev. **D75**, 115005 (2007), arXiv:hep-ph/0703097 [HEP-PH].
- [10] S. P. Martin, Phys.Rev. **D77**, 075002 (2008), arXiv:0801.0237 [hep-ph].
- [11] T. J. LeCompte and S. P. Martin, Phys.Rev. **D85**, 035023 (2012), arXiv:1111.6897 [hep-ph].
- [12] G. Belanger, M. Heikinheimo, and V. Sanz, JHEP **1208**, 151 (2012), arXiv:1205.1463 [hep-ph].
- [13] K. Rolbiecki and K. Sakurai, JHEP **1210**, 071 (2012), arXiv:1206.6767 [hep-ph].
- [14] D. Curtin, P. Meade, and P.-J. Tien, Phys.Rev. **D90**, 115012 (2014), arXiv:1406.0848 [hep-ph].
- [15] J. S. Kim, K. Rolbiecki, K. Sakurai, and J. Tattersall, JHEP **1412**, 010 (2014), arXiv:1406.0858 [hep-ph].
- [16] M. Czakon, A. Mitov, M. Papucci, J. T. Ruderman, and A. Weiler, Phys.Rev.Lett. **113**, 201803 (2014), arXiv:1407.1043 [hep-ph].
- [17] V. Khachatryan *et al.* (CMS), Phys.Lett. **B743**, 503 (2015), arXiv:1411.7255 [hep-ex].
- [18] K. Rolbiecki and J. Tattersall, (2015), arXiv:1505.05523 [hep-ph].
- [19] H. An and L.-T. Wang, (2015), arXiv:1506.00653 [hep-ph].
- [20] Z. Chacko, H.-S. Goh, and R. Harnik, Phys.Rev.Lett. **96**, 231802 (2006), arXiv:hep-ph/0506256 [hep-ph].
- [21] G. Burdman, Z. Chacko, H.-S. Goh, and R. Harnik, JHEP **0702**, 009 (2007), arXiv:hep-ph/0609152 [hep-ph].
- [22] H. Cai, H.-C. Cheng, and J. Terning, JHEP **0905**, 045 (2009), arXiv:0812.0843 [hep-ph].
- [23] R. Barbieri, T. Gregoire, and L. J. Hall, (2005), arXiv:hep-ph/0509242 [hep-ph].
- [24] Z. Chacko, Y. Nomura, M. Papucci, and G. Perez, JHEP **01**, 126 (2006), arXiv:hep-ph/0510273 [hep-ph].
- [25] I. G. Garca, R. Lasenby, and J. March-Russell, (2015), arXiv:1505.07109 [hep-ph].
- [26] N. Craig and A. Katz, (2015), arXiv:1505.07113 [hep-ph].
- [27] I. G. Garca, R. Lasenby, and J. March-Russell, (2015), arXiv:1505.07410 [hep-ph].
- [28] M. Farina, (2015), arXiv:1506.03520 [hep-ph].
- [29] P. Schwaller, (2015), arXiv:1504.07263 [hep-ph].
- [30] D. Poland and J. Thaler, JHEP **0811**, 083 (2008), arXiv:0808.1290 [hep-ph].
- [31] B. Batell and M. McCullough, (2015), arXiv:1504.04016 [hep-ph].
- [32] C. Csaki, M. Geller, O. Telem, and A. Weiler, (2015), arXiv:1512.03427 [hep-ph].
- [33] N. Craig, S. Knapen, and P. Longhi, Phys.Rev.Lett. **114**, 061803 (2015), arXiv:1410.6808 [hep-ph].
- [34] N. Craig, S. Knapen, and P. Longhi, JHEP **1503**, 106 (2015), arXiv:1411.7393 [hep-ph].
- [35] P. Batra and Z. Chacko, Phys.Rev. **D79**, 095012 (2009), arXiv:0811.0394 [hep-ph].
- [36] R. Barbieri, D. Greco, R. Rattazzi, and A. Wulzer, (2015), arXiv:1501.07803 [hep-ph].
- [37] M. Low, A. Tesi, and L.-T. Wang, Phys.Rev. **D91**, 095012 (2015), arXiv:1501.07890 [hep-ph].
- [38] M. Geller and O. Telem, Phys.Rev.Lett. **114**, 191801 (2015), arXiv:1411.2974 [hep-ph].
- [39] N. Craig and K. Howe, JHEP **1403**, 140 (2014), arXiv:1312.1341 [hep-ph].
- [40] N. Craig and H. K. Lou, JHEP **1412**, 184 (2014), arXiv:1406.4880 [hep-ph].
- [41] S. Chang, L. J. Hall, and N. Weiner, Phys. Rev. **D75**, 035009 (2007), arXiv:hep-ph/0604076 [hep-ph].
- [42] D. Curtin and P. Saraswat, (2015), arXiv:1509.04284 [hep-ph].
- [43] H.-C. Cheng, S. Jung, E. Salvioni, and Y. Tsai, (2015), arXiv:1512.02647 [hep-ph].
- [44] M. J. Strassler and K. M. Zurek, Phys.Lett. **B651**, 374 (2007), arXiv:hep-ph/0604261 [hep-ph].
- [45] M. J. Strassler and K. M. Zurek, Phys.Lett. **B661**, 263 (2008), arXiv:hep-ph/0605193 [hep-ph].
- [46] M. J. Strassler, (2006), arXiv:hep-ph/0607160 [hep-ph].
- [47] T. Han, Z. Si, K. M. Zurek, and M. J. Strassler, JHEP **0807**, 008 (2008), arXiv:0712.2041 [hep-ph].
- [48] C. J. Morningstar and M. J. Peardon, Phys.Rev. **D60**, 034509 (1999), arXiv:hep-lat/9901004 [hep-lat].
- [49] J. E. Lukneevich, D. Melnikov, and M. J. Strassler, JHEP **0907**, 055 (2009), arXiv:0903.0883 [hep-ph].
- [50] D. Curtin and C. B. Verhaaren, (2015), arXiv:1506.06141 [hep-ph].
- [51] C. Csaki, E. Kuflik, S. Lombardo, and O. Slone, Phys. Rev. **D92**, 073008 (2015), arXiv:1508.01522 [hep-ph].
- [52] G. Burdman, Z. Chacko, R. Harnik, L. de Lima, and C. B. Verhaaren, Phys.Rev. **D91**, 055007 (2015), arXiv:1411.3310 [hep-ph].
- [53] D. Curtin, R. Essig, S. Gori, P. Jaiswal, A. Katz, *et al.*, Phys.Rev. **D90**, 075004 (2014), arXiv:1312.4992 [hep-ph].
- [54] L. B. Okun, JETP Lett. **31**, 144 (1980), [Pisma Zh. Eksp. Teor. Fiz.31,156(1979)].
- [55] L. B. Okun, Nucl. Phys. **B173**, 1 (1980).
- [56] J. Kang and M. A. Luty, JHEP **0911**, 065 (2009), arXiv:0805.4642 [hep-ph].
- [57] G. Burdman, Z. Chacko, H.-S. Goh, R. Harnik, and C. A. Krenke, Phys.Rev. **D78**, 075028 (2008), arXiv:0805.4667 [hep-ph].
- [58] K. Cheung, W.-Y. Keung, and T.-C. Yuan, Nucl.Phys. **B811**, 274 (2009), arXiv:0810.1524 [hep-ph].
- [59] R. Harnik and T. Wizansky, Phys.Rev. **D80**, 075015 (2009), arXiv:0810.3948 [hep-ph].
- [60] R. Harnik, G. D. Kribs, and A. Martin, Phys.Rev. **D84**, 035029 (2011), arXiv:1106.2569 [hep-ph].
- [61] R. Fok and G. D. Kribs, Phys. Rev. **D84**, 035001 (2011), arXiv:1106.3101 [hep-ph].
- [62] Z. Chacko, D. Curtin, and C. B. Verhaaren, (2016).
- [63] Z. Chacko, D. Curtin, and C. B. Verhaaren, (2016).

- [64] T. Cohen, N. Craig, H. K. Lou, and D. Pinner, (2015), arXiv:1508.05396 [hep-ph].
- [65] G. Burdman and R. T. D'Agnolo, (2015), arXiv:1512.00040 [hep-ph].
- [66] S. Dimopoulos, K. Howe, and J. March-Russell, Phys. Rev. Lett. **113**, 111802 (2014), arXiv:1404.7554 [hep-ph].
- [67] I. G. Garcia, K. Howe, and J. March-Russell, (2015), arXiv:1510.07045 [hep-ph].
- [68] J. Alwall, M. Herquet, F. Maltoni, O. Mattelaer, and T. Stelzer, JHEP **1106**, 128 (2011), arXiv:1106.0522 [hep-ph].
- [69] B. Batell and S. Jung, (2015), arXiv:1504.01740 [hep-ph].
- [70] L. Clavelli, Phys. Lett. **B85**, 111 (1979).
- [71] J. del Peso, L. Labarga, and F. Barreiro, Z. Phys. **C43**, 287 (1989).
- [72] R. K. Ellis, W. J. Stirling, and B. R. Webber, Camb. Monogr. Part. Phys. Nucl. Phys. Cosmol. **8**, 1 (1996).
- [73] J. Juknevich, *Phenomenology of pure-gauge hidden valleys at Hadron colliders*, Ph.D. thesis, Rutgers U., Piscataway.
- [74] V. D. Barger, E. W. N. Glover, K. Hikasa, W.-Y. Keung, M. G. Olsson, C. J. Suchyta, III, and X. R. Tata, Phys. Rev. **D35**, 3366 (1987), [Erratum: Phys. Rev.D38,1632(1988)].
- [75] P. Schwaller, D. Stolarski, and A. Weiler, JHEP **1505**, 059 (2015), arXiv:1502.05409 [hep-ph].